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On The BV Formulation Of Boundary Superstring Field Theory

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We propose a Batalin-Vilkovisky (BV) formulation of boundary superstring field theory. The superstring field action is defined in terms of a closed one-form in the space of couplings, and we compute it explicitly for exactly solvable tachyon perturbations. We also argue that the superstring field action defined in this way is the partition function on the disc, in accord with a previous proposal.

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1. Introduction

The problem of open string tachyon condensation has been studied intensively in the last two years. String field theory has shown to be a very fruitful framework to address this problem, and it has provided strong evidence for Sen's conjectures [1]. There are currently two different approaches to (bosonic) string field theory: the cubic string field theory formulated by Witten [2], and the so-called boundary string field theory (BSFT) of Witten and Shatashvili [3][4][5]. In the cubic theory, the approach to tachyon condensation has relied so far in an approximation scheme called level truncation [6][7][1]. On the other hand, BSFT provides exact results on the tachyon potential [9][10] and on D-brane tensions [10].

The bosonic BSFT has a nice geometric formulation in terms of the Batalin-Vilkovisky (BV) formalism [3]. In this formulation, the string field action S is defined up to a constant by a locally exact one-form dS in the space of couplings for boundary perturbations. To define dS, one needs a closed two-form ω of ghost number -1, together with a vector field V of ghost number 1 which is a symmetry of ω (*i.e.*, $\mathcal{L}_V \omega = 0$). When this is the case, the string field action is defined by

$$dS = \iota_V \, \omega, \tag{1.1}$$

which is indeed closed (since $d\omega = \mathcal{L}_V \omega = 0$), and then locally exact.

In BSFT, the vector field V and two-form are defined as follows (when ghosts and matter are decoupled). Take a worldsheet action with the structure

$$I = I_{\text{bulk}} + \int_{\partial\Sigma} \mathcal{V}, \tag{1.2}$$

where I_{bulk} is the standard closed string background for matter and ghost, and $\mathcal{V}$ is a generic matter operator. Consider the ghost-number 1 operator $\mathcal{O} = c\mathcal{V}$. We then define the two-form ω as follows:

$$\omega(\mathcal{O}_1, \mathcal{O}_2) = \frac{1}{2} \oint d\tau \oint d\tau' \langle \mathcal{O}_1(\tau) \mathcal{O}_2(\tau') \rangle, \tag{1.3}$$

where the brackets denote correlation functions in the theory with worldsheet action (1.2). Since the ghost number of the vacuum is -3, ω has indeed ghost number -1. The vector

[1] A different approach has been proposed recently in [8].

field V is defined by the BRST operator $Q_{\rm BRST}$ in the bulk theory, which has ghost number 1, and (1.1) gives:

$$dS = \frac{1}{2} \oint d\tau \oint d\tau' \langle d\mathcal{O}(\tau) \{ Q_{\rm BRST}, \mathcal{O}(\tau') \} \rangle. \tag{1.4}$$

The fact that $d\omega = \mathcal{L}_V \omega = 0$ are consequences of rotation invariance, and invariance under b_{-1} and $Q_{\rm BRST}$. The equation (1.4) defines the spacetime string field action in BSFT. Starting from (1.4), one can show that S is given by [4][5]:

$$S = \big(\beta^k \frac{\partial}{\partial x^k} + 1 \big) Z, \tag{1.5}$$

where the x^k are couplings for boundary operators, and β^k are the corresponding β-functions.

In the superstring case, the level truncation scheme has been also developed [11] in the superstring field theory formulated by Berkovits [12], as well as in some other models [13], providing again some important evidence for Sen's conjectures. In [14], the problem of constructing a boundary superstring field theory (BSSFT) was addressed. Based on some previous works on the sigma model approach to superstring field theory [15], it was proposed that the spacetime superstring field action is given by the (perturbed) partition function on the disc. This is in contrast with the bosonic case (1.5), where the string field action differs from the partition function in a term involving the β-functions.

In this paper we make a few steps towards a geometric construction (a la BV) of BSSFT. We restrict ourselves to the NS sector of the open superstring, and to perturbations where the ghosts and the matter are decoupled. In section 2 we make a proposal for dS which is a simple generalization of (1.4) to the superstring case. In section 3, we make some detailed computations for exactly solvable tachyon perturbations, and find that $S = Z$. In section 4, we give a simple and general argument to show that $S = Z$. This gives further support to the proposal made in [14] from a different, more geometric point of view. We end in section 5 with some open problems.

2. Proposal for BSSFT

In order to find a BV formulation of BSSFT, which generalizes the construction in [3], we have to construct a closed one-form in the space of coupling constants. We will restrict ourselves to the NS sector of the open superstring. We take as the perturbed action on the disc,

$$I = I_{\rm bulk} + I_{\Gamma} + \int_{\partial\Sigma} G_{-1/2} \mathcal{O}. \tag{2.1}$$

In this equation, $I_{\rm bulk}$ is the RNS bulk action including the bc and the $\beta\gamma$ (super)ghosts. I_Γ is the action for the auxiliary fermionic superfield living in the boundary [16][10]:

$$I_\Gamma = \oint \frac{d\tau}{2\pi} d\theta \Gamma D \Gamma, \tag{2.2}$$

where

$$\Gamma = \mu + \theta F \tag{2.3}$$

and $D = \partial_\theta + \theta\partial_\tau$. This auxiliary superfield is needed in order to incorporate the GSO(−) states. Finally, $\mathcal{O}$ is the lowest component of a worldsheet superfield:

$$\Psi = \mathcal{O} + \theta G_{-1/2}\mathcal{O}, \tag{2.4}$$

i.e. $G_{-1/2}\mathcal{O}$ is a NS matter operator in the 0-picture.

A natural proposal for the superstring field action is:

$$dS = \frac{1}{8}\oint \frac{d\tau d\tau'}{(2\pi)^2} \langle c(\tau) d\mathcal{V}^{(-1)}(\tau)\{Q_{\rm BRST}, c(\tau')\mathcal{V}^{(-1)}(\tau')\}\rangle, \tag{2.5}$$

where

$$\mathcal{V}^{(-1)} = {\rm e}^{-\phi}\mathcal{O} \tag{2.6}$$

is the usual vertex in the (-1)-picture, and the correlation function is computed in the theory with action (2.1). The operator $Q_{\rm BRST}$ is the BRST operator of the NSR string [17], and we assume that it does not act on the auxiliary field Γ. This is a natural assumption, since Γ comes from Chan-Paton degrees of freedom, and lives in the boundary, while $Q_{\rm BRST}$ comes from the bulk CFT. Notice that we have the appropriate number of (super)ghost insertions to soak up all the zero modes, and that we are working in the "small" Hilbert space, as in the formulation given in [18].

The above proposal is a natural generalization of (1.4). The vector field in the space of couplings is again given by the BRST operator $Q_{\rm BRST}$, and the two-form ω is a natural generalization of (1.3):

$$\omega(\mathcal{O}_1, \mathcal{O}_2) = \frac{1}{8}\oint \frac{d\tau d\tau'}{(2\pi)^2} \langle c(\tau){\rm e}^{-\phi(\tau)}\mathcal{O}_1(\tau) c(\tau'){\rm e}^{-\phi(\tau')}\mathcal{O}_2(\tau')\rangle, \tag{2.7}$$

after taking into account the superconformal ghosts. Since ${\rm e}^{-\phi}$ has zero ghost number, ω has again ghost number -1. It follows from the arguments in [3] that dS is a closed one-form.

3. Some explicit computations

In order to illustrate the above proposal, we will compute dS explicitly when the boundary perturbation is a constant or a linear tachyon field. As it has been shown it [14], in these cases the theory is exactly solvable (since it is a free theory). The superfield describing a tachyon profile is [16][14]:

$$\Psi = \Gamma\, T(\mathbf{X}), \tag{3.1}$$

therefore

$$\mathcal{O} = \mu\, T(X), \tag{3.2}$$

and

$$G_{-1/2}\mathcal{O} = F\, T(X) + \psi^\rho \mu \partial_\rho T. \tag{3.3}$$

The action for the auxiliary superfield and the boundary perturbation add up to [14]:

$$\oint \frac{d\tau}{2\pi}\big(F^2 + \dot\mu\mu + FT(X) + \psi^\rho\mu\partial_\rho T\big). \tag{3.4}$$

3.1. Constant tachyon

The computation for the constant tachyon is particularly simple. First, we have

$$\{Q_{\rm BRST}, c\mu T(X)\} = \frac{1}{2} c\partial c \mathrm{e}^{-\phi}\mu T(X) + 2c\partial c\mathrm{e}^{-\phi}\mu\frac{\partial^2 T}{\partial X^2}. \tag{3.5}$$

Since T is a constant, we find that

$$dS = -\frac{1}{16}\oint \frac{d\tau d\tau'}{(2\pi)^2}\langle c\partial c(\tau)c(\tau')\rangle\langle \mathrm{e}^{-\phi}(\tau)\mathrm{e}^{-\phi}(\tau')\rangle\langle\mu(\tau)\mu(\tau')\rangle T\, dT \exp\big\{-\frac{1}{4}T^2\big\}. \tag{3.6}$$

In this equation, the correlators are evaluated in the decoupled bc, ϕ and μ theories, and we have integrated out the auxiliary field F as in [10]. Using now the correlation functions,

$$\begin{aligned}\langle c\partial c(\tau)c(\tau')\rangle &= -\,4\,\sin^2\left(\frac{\tau-\tau'}{2}\right),\\ \langle \mathrm{e}^{-\phi}(\tau)\mathrm{e}^{-\phi}(\tau')\rangle &= -\,\frac{2}{\sin\big(\frac{\tau-\tau'}{2}\big)},\\ \langle\mu(\tau)\mu(\tau')\rangle &= \frac{\pi}{2}\epsilon(\tau-\tau').\end{aligned} \tag{3.7}$$

one obtains:

$$dS = d\big(\mathrm{e}^{-\frac{1}{4}T^2}\big). \tag{3.8}$$

Integrating (3.8), we find that $S = Z$ (we have not included in (3.8) a global normalization proportional to the volume of the D9-brane).

3.2. Linear tachyon

The computation of dS for a linear tachyon

$$T(X) = uX \tag{3.9}$$

is more involved. The definition gives,

$$dS = -\frac{1}{16}\int \frac{d\tau d\tau'}{(2\pi)^2}\langle c\partial c(\tau)c(\tau')\rangle\langle \mathrm{e}^{-\phi}(\tau)\mathrm{e}^{-\phi}(\tau')\rangle\langle\mu(\tau)\mu(\tau')\rangle udu\langle X(\tau)X(\tau')\rangle. \tag{3.10}$$

Now the action (3.4) for μ is a Gaussian with a linear term. We then shift,

$$\tilde{\mu} = \mu + \frac{1}{2}\frac{1}{\partial_\tau}\psi^\rho\partial_\rho T, \tag{3.11}$$

where

$$\frac{1}{\partial_\tau}f(\tau) = \frac{1}{2}\int d\tau'\epsilon(\tau-\tau')f(\tau'). \tag{3.12}$$

Using all this, (3.10) becomes the sum of two terms. The first one is simply,

$$dS_1 = -\frac{\pi}{8}dyZ(y)\int\frac{d\tau d\tau'}{(2\pi)^2}\sin\left(\frac{\tau-\tau'}{2}\right)\epsilon(\tau-\tau')G_B(\tau-\tau',y), \tag{3.13}$$

where $y = u^2$, $Z(y)$ is the partition function, and $G_B(\tau-\tau',y)$ is the perturbed bosonic propagator [4],

$$G_B(\tau-\tau',y) = 2\sum_{k\in\mathbf{Z}}\frac{1}{|k|+y}\mathrm{e}^{ik(\tau-\tau')}. \tag{3.14}$$

The second piece is:

$$dS_2 = -\frac{1}{64}ydyZ(y)\int\frac{d\tau d\tau'}{(2\pi)^2}d\sigma d\sigma'\sin\left(\frac{\tau-\tau'}{2}\right)\epsilon(\tau-\sigma)\epsilon(\tau'-\sigma')G_B(\tau-\tau',y)G_F(\sigma-\sigma',y), \tag{3.15}$$

where [10]

$$G_F(\tau-\tau',y) = 2i\sum_{k\in\mathbf{Z}+\frac{1}{2}}\frac{k}{|k|+y}\mathrm{e}^{ik(\tau-\tau')} \tag{3.16}$$

is the perturbed fermionic propagator. Using the explicit representation of the step function as

$$\epsilon(\tau-\tau') = \frac{2}{\pi}\sum_{k\in\mathbf{Z}+\frac{1}{2}>0}\frac{\sin k(\tau-\tau')}{r}, \tag{3.17}$$

we find for the first term

$$dS_1 = \frac{1}{4} dy Z(y) \Big\{ - \sum_{n=0}^{\infty} \frac{1}{(n+y)(n+1/2)} + \sum_{n=1}^{\infty} \frac{1}{(n+y)(n-1/2)} \Big\}, \tag{3.18}$$

and for the second term,

$$\begin{aligned} dS_2 =& \frac{1}{4} y dy Z(y) \Big\{ \sum_{n=0}^{\infty} \frac{1}{(n+y)(n+y+1/2)(n+1/2)} \\ & - \sum_{n=1}^{\infty} \frac{1}{(n+y)(n+y-1/2)(n-1/2)} \Big\}. \end{aligned} \tag{3.19}$$

Therefore, the one-form (2.5) is in this case:

$$dS = -\frac{1}{4} dy Z(y) \Big\{ \sum_{n=0}^{\infty} \frac{1}{(n+y)(n+y+1/2)} - \sum_{n=1}^{\infty} \frac{1}{(n+y)(n+y-1/2)} \Big\}. \tag{3.20}$$

Notice that all the infinite sums involved here are convergent. We can explicitly evaluate them by decomposing in simple fractions and using the formula

$$\psi(x) - \psi(y) = \sum_{n=0}^{\infty} \Big(\frac{1}{y+n} - \frac{1}{x+n} \Big), \tag{3.21}$$

where $\psi(x)$ is the logarithmic derivative of the Γ function, as well as the doubling formula

$$\psi(2x) = \log 2 + \frac{1}{2} \big(\psi(x) + \psi(x+1/2) \big). \tag{3.22}$$

The infinite sums in (3.20) add up to:

$$2 \big\{ -4 \log 2 - \frac{1}{y} - 4 \big(\psi(y) - \psi(2y) \big) \big\}. \tag{3.23}$$

Using the explicit results of [14], it is easy to see that this is nothing but the correlator

$$\langle X^2 + \psi \frac{1}{\partial_\tau} \psi \rangle \tag{3.24}$$

regularized with a point-splitting procedure which preserves supersymmetry. We then find:

$$dS = -\frac{1}{4} \langle X^2 + \psi \frac{1}{\partial_\tau} \psi \rangle Z(y) dy, \tag{3.25}$$

On the other hand, it was shown in [14] that

$$\frac{d \log Z}{dy} = -\frac{1}{4} \langle X^2 + \psi \frac{1}{\partial_\tau} \psi \rangle. \tag{3.26}$$

Therefore, we can integrate dS to obtain

$$S = Z(y), \tag{3.27}$$

up to an additive constant. It is interesting to note that the convergent sums involved in dS give automatically the supersymmetric regularization of the propagators proposed in [14].

4. $S = Z$

The explicit computations of the previous section suggest that the action defined in (2.5) is in fact the partition function on the disc, *i.e.* $S = Z$. In this section we give a general argument showing that this is in fact the case, at least in the case in which matter and ghosts are decoupled. To do this, we adapt the argument given in [4]. The idea is to consider two decoupled subsystems, with partition functions Z_1 and Z_2, in such a way that the combined partition function is $Z = Z_1 Z_2$. Let us now evaluate dS in the case in which we turn on a series of operators in the GSO($-$) sector. If $\mathcal{O}_i$ denotes a basis of operators for the first subsystem, and $\widetilde{\mathcal{O}}_j$ a basis of operators in the second subsystem, then the general operator $\mathcal{O}$ has the form:

$$\mathcal{O} = \sum_i x^i \mathcal{O}_i + \sum_j y^j \widetilde{\mathcal{O}}_j. \tag{4.1}$$

In this equation, x^i and y^j are coupling constants for the first and second subsystems, respectively, and the operators in the GSO($-$) sector have the form,

$$\mathcal{O}_i = \mu_1 \mathcal{B}_i, \ \widetilde{\mathcal{O}}_j = \mu_2 \widetilde{\mathcal{B}}_j, \tag{4.2}$$

where μ_1, μ_2 are the boundary fermions for the first and second subystems. and $\mathcal{B}_i$, $\widetilde{\mathcal{B}}_j$ are bosonic matter operators for the first and second subsystems. The key fact is that, since the two subsystems are decoupled,

$$\langle \mu_1(\tau)\, \mu_2(\tau') \rangle = 0. \tag{4.3}$$

By assumption, the BRST charge does not act on the boundary fermions, and after evaluating the ghost correlation functions, we find:

$$\begin{aligned} dS = \oint d\tau d\tau' \langle d\mathcal{O}(\tau) \Big[\big(\sum_i A^i(x, \tau-\tau') \mathcal{B}_i(\tau') \big) \mu_1(\tau') \\ + \big(\sum_j \widetilde{A}^j(y, \tau-\tau') \widetilde{\mathcal{B}}_j(\tau') \big) \mu_2(\tau') \Big] \rangle, \end{aligned} \tag{4.4}$$

where $A^i(x, \tau-\tau')$, $\widetilde{A}^j(y, \tau-\tau')$ are some functions for the first and second subsystem, respectively, whose detailed structure will not be relevant for the argument. After evaluating the correlation functions and performing the integrals, we find, using (4.3), that dS can be written as

$$dS = \big(\sum_i dx^i a_i(x) \big) Z_2 + Z_1 \big(\sum_j dy^j \widetilde{a}_j(y) \big), \tag{4.5}$$

where again $a_i(x)$, $\widetilde{a}_j(y)$ are some functions of the couplings for the first and second subsystems. Notice that, in contrast to the situation analyzed in [4], the presence of the decoupled boundary fermions implies that correlation functions of the form $\langle \mathcal{O}_i \widetilde{\mathcal{O}}_j \rangle$ vanish and do not appear in the final expression for dS.

We now use the key fact that dS is a closed one-form, *i.e.* $d^2 S = 0$. Setting to zero the coefficient of the $dx^i \wedge dy^j$ terms gives

$$-\big(\sum_i dx^i a_i(x)\big) dZ_2 + dZ_1 \big(\sum_j dy^j \widetilde{a}_j(y)\big) = 0. \tag{4.6}$$

Since $dZ_{1,2}$ only depends on x (respectively, y), we find that

$$g dZ_1 = \sum_i dx^i a_i(x), \quad g dZ_2 = \sum_j dy^j \widetilde{a}_j(y), \tag{4.7}$$

for some constant g. This fixes the unknown functions $a_i(x)$, $\widetilde{a}_j(y)$ in terms of the partition functions of the subsystems. In particular, one finds that, up to an additive constant,

$$S = g\, Z_1\, Z_2. \tag{4.8}$$

The constant g can be fixed by computing S for some particular couplings, as we did before. We therefore find $g = 1$ and finally

$$S = Z. \tag{4.9}$$

This argument can be extended to perturbations in the GSO(+) sector. The key ingredient is that the operator $\mathcal{O}$ is again fermionic. Since the two subsystems are decoupled, correlation functions of the form $\langle \mathcal{O}_i \widetilde{\mathcal{O}}_j \rangle$ vanish, and the above argument goes through. This shows that $S = Z$ for operators in the NS sector, when matter and ghosts are decoupled.

5. Open problems

In this note we have made a proposal for a BV formulation of BSSFT, and we have shown that the spacetime superstring field action turns out to be the partition function on the disc, in agreement with [14]. There are many problems related to this formulation that have not been addressed. Some of these problems are the following:

1) As noticed in [5][10], the BV formulation of BSFT is not covariant. The covariant definition of the string field action is rather,

$$\frac{\partial S}{\partial x^i} = \beta^j \mathcal{G}_{ij}, \tag{5.1}$$

where $\mathcal{G}_{ij}$ is a positive-definite metric. It would be very interesting to see if our definition admits such a covariant extension, involving a metric which is manifestly positive-definite. This would immediately show that the superstring field action is monotonically decreasing along RG flows, as in the bosonic case.

2) Here we have concentrated on the non-projected open NSR superstring, which describes the non-BPS D9-brane of type IIA theory. It would be interesting to extend our considerations to the D-$\overline{\mathrm{D}}$ system, which has been analyzed from the point of view of BSFT in [19][20].

3) An important open problem is the inclusion of the R sector in this formulation of BSSFT. This is a problematic issue in the existing formulations of superstring field theory, and BSSFT may turn out to be a better framework.

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